DESY 01–006
CERN–TH/2001–012
TTP01–06
hep–ph/0101276
January 2001

General Analysis of New Physics in $B \to J/\psi K$

Robert Fleischer*
Deutsches Elektronen-Synchrotron DESY, Notkestraße 85, D–22607 Hamburg, Germany

Thomas Mannel†
CERN Theory Division, CH–1211 Geneva 23, Switzerland
and
Institut für Theoretische Teilchenphysik, Universität Karlsruhe,
D–76128 Karlsruhe, Germany

Abstract

We present a model-independent parametrization of the $B^\pm \to J/\psi K^\pm$, $B_d \to J/\psi K_S$ decay amplitudes by taking into account the constraints that are implied by the isospin symmetry of strong interactions. Employing estimates borrowed from effective field theory, we explore the impact of physics beyond the Standard Model and introduce – in addition to the usual mixing-induced CP asymmetry $\mathcal{A}_{CP}^{mix}$ in $B_d \to J/\psi K_S$ – a set of three observables, allowing a general analysis of possible new-physics effects in the $B \to J/\psi K$ system. Imposing a dynamical hierarchy of amplitudes, we argue that one of these observables may already be accessible at the first-generation B-factories, whereas the remaining ones will probably be left for second-generation B experiments. However, in the presence of large rescattering effects, all three new-physics observables may be sizeable. We also emphasize that a small value of $\mathcal{A}_{CP}^{mix}$ could be due to new-physics effects arising at the $B \to J/\psi K$ decay-amplitude level. In order to establish such a scenario, the observables introduced in this paper play a key role.

*E-mail: Robert.Fleischer@desy.de
†E-mail: Thomas.Mannel@cern.ch

1 Introduction

At present, we are at the beginning of the B-factory era in particle physics, which will provide valuable insights into CP violation and various tests of the Kobayashi–Maskawa picture of this phenomenon [1]. Among the most interesting B-decay channels is the "gold-plated" mode $B_d \to J/\psi K_S$ [2], which allows the determination of the angle β of the unitarity triangle of the Cabibbo–Kobayashi–Maskawa (CKM) matrix [3]. In the summer of 2000 – after important steps at LEP [4] and by the CDF collaboration [5] – the first results on CP-violating effects in $B_d \to J/\psi K_S$ were reported by the BaBar [6] and Belle [7] collaborations, which already led to some excitement in the B-physics community [8].

In this paper, we consider the neutral mode $B_d \to J/\psi K_S$ together with its charged counterpart $B^\pm \to J/\psi K^\pm$. Making use of the isospin symmetry of strong interactions, we derive a model-independent parametrization of the corresponding decay amplitudes. After a careful analysis of the Standard-Model contributions, we include possible new-physics amplitudes and introduce – in addition to the usual "mixing-induced" CP asymmetry in $B_d \to J/\psi K_S$ – a set of three observables, allowing a general analysis of new-physics effects in the $B \to J/\psi K$ system; the generic size of these effects is estimated with the help of arguments borrowed from effective field theory. The four observables provided by $B \to J/\psi K$ decays are affected by physics beyond the Standard Model in two different ways: through B^0_d–$\overline{B^0_d}$ mixing, and new-physics contributions to the $B \to J/\psi K$ decay amplitudes. Usually, B^0_d–$\overline{B^0_d}$ mixing is considered as the preferred mechanism for new physics to manifest itself in the mixing-induced $B_d \to J/\psi K_S$ CP asymmetry [9]. Here we focus on new-physics effects arising at the $B \to J/\psi K$ decay-amplitude level [10]; we emphasize that the extraction of the CKM angle β from mixing-induced CP violation in $B_d \to J/\psi K_S$ may be significantly disturbed by such effects. As is well known, the value of "β" itself, i.e. the CP-violating weak B^0_d–$\overline{B^0_d}$ mixing phase, may deviate strongly from the Standard-Model expectation because of new-physics contributions to B^0_d–$\overline{B^0_d}$ mixing.

The three observables introduced in this paper allow us to search for "smoking-gun" signals of new-physics contributions to the $B \to J/\psi K$ decay amplitudes, which would also play an important role for mixing-induced CP violation in $B_d \to J/\psi K_S$. Employing an isospin decomposition and imposing a hierarchy of amplitudes, we argue that one of these observables may be sizeable and could already be accessible at the first-generation B-factories. On the other hand, the remaining two observables are expected to be dynamically suppressed, but may be within reach of the second-generation B experiments, BTeV and LHCb. In the presence of large rescattering effects, all three observables could be of the same order of magnitude. However, we do not consider this as a very likely scenario and note that also the "QCD factorization" approach is not in favour of such large rescattering processes [11].

The outline of this paper is as follows: in Section 2, we investigate the isospin structure of the $B^\pm \to J/\psi K^\pm$, $B_d \to J/\psi K_S$ decay amplitudes, and parametrize them in a model-independent manner; a particular emphasis is given to the Standard-Model case. In Section 3, we then have a closer look at the impact of new physics. To this end, we make use of dimensional estimates following from effective field theory, and introduce a plausible dynamical hierarchy of amplitudes. The set of observables to search for "smoking-gun" signals of new-physics contributions to the $B \to J/\psi K$ decay amplitudes is defined in Section 4, and is discussed in further detail in Section 5. In Section 6, our main points are summarized.

2 Phenomenology of $B \to J/\psi K$ Decays

The most general discussion of the $B^{\pm} \to J/\psi K^{\pm}$, $B_d \to J/\psi K_{\rm S}$ system can be performed in terms of an isospin decomposition. Here the corresponding initial and final states are grouped in the following isodoublets:

$$\begin{pmatrix} |1/2;+1/2\rangle \\ |1/2;-1/2\rangle \end{pmatrix}: \quad \underbrace{\begin{pmatrix} |B^+\rangle \\ |B_d^0\rangle \end{pmatrix}, \quad \begin{pmatrix} |\overline{B_d^0}\rangle \\ -|B^-\rangle \end{pmatrix}}_{\mathcal{CP}}, \quad \underbrace{\begin{pmatrix} |J/\psi K^+\rangle \\ |J/\psi K^0\rangle \end{pmatrix}, \quad \begin{pmatrix} |J/\psi \overline{K^0}\rangle \\ -|J/\psi K^-\rangle \end{pmatrix}}_{\mathcal{CP}}, \tag{1}$$

which are related by CP conjugation. As usual, the $K_{\rm S}$ state in $B_d \to J/\psi K_{\rm S}$ is the superposition of the two neutral kaon states $|K^0\rangle$ and $|\overline{K^0}\rangle$ corresponding to CP eigenvalue $+1$. The decays $B^+ \to J/\psi K^+$ and $B_d^0 \to J/\psi K^0$ are described by an effective low-energy Hamiltonian of the following structure:

$$\mathcal{H}_{\rm eff} = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs}V_{cb}^* \left(\mathcal{Q}_{\rm CC}^c - \mathcal{Q}_{\rm QCD}^{\rm pen} - \mathcal{Q}_{\rm EW}^{\rm pen} \right) + V_{us}V_{ub}^* \left(\mathcal{Q}_{\rm CC}^u - \mathcal{Q}_{\rm QCD}^{\rm pen} - \mathcal{Q}_{\rm EW}^{\rm pen} \right) \right], \tag{2}$$

where the $\mathcal{Q}$ are linear combinations of perturbatively calculable Wilson coefficient functions and four-quark operators, consisting of current–current (CC), QCD penguin and electroweak (EW) penguin operators. For an explicit list of operators and a detailed discussion of the derivation of (2), the reader is referred to [12]. For the following considerations, the flavour structure of these operators plays a key role:

$$\mathcal{Q}_{\rm CC}^c \sim (\overline{c}c)(\overline{b}s), \quad \mathcal{Q}_{\rm CC}^u \sim (\overline{u}u)(\overline{b}s), \tag{3}$$

$$\mathcal{Q}_{\rm QCD}^{\rm pen} \sim \left[(\overline{c}c) + \{(\overline{u}u) + (\overline{d}d)\} + (\overline{s}s) \right] (\overline{b}s), \tag{4}$$

$$\mathcal{Q}_{\rm EW}^{\rm pen} \sim \frac{1}{3} \left[2(\overline{c}c) + \{2(\overline{u}u) - (\overline{d}d)\} - (\overline{s}s) \right] (\overline{b}s), \tag{5}$$

where the factors of $+2/3$ and $-1/3$ in (5) are due to electrical quark charges. In (3)–(5), we have suppressed all colour and spin indices. Since

$$(\overline{u}u) = \frac{1}{2} \underbrace{\left(\overline{u}u + \overline{d}d\right)}_{I=0} + \frac{1}{2} \underbrace{\left(\overline{u}u - \overline{d}d\right)}_{I=1}, \quad 2(\overline{u}u) - (\overline{d}d) = \frac{1}{2} \underbrace{\left(\overline{u}u + \overline{d}d\right)}_{I=0} + \frac{3}{2} \underbrace{\left(\overline{u}u - \overline{d}d\right)}_{I=1}, \tag{6}$$

we conclude that the Hamiltonian (2) is a combination of isospin $I=0$ and $I=1$ pieces:

$$\mathcal{H}_{\rm eff} = \mathcal{H}_{\rm eff}^{I=0} + \mathcal{H}_{\rm eff}^{I=1}, \tag{7}$$

where $\mathcal{H}_{\rm eff}^{I=0}$ receives contributions from all of the operators listed in (3)–(5), whereas $\mathcal{H}_{\rm eff}^{I=1}$ is only due to $\mathcal{Q}_{\rm CC}^u$ and $\mathcal{Q}_{\rm EW}^{\rm pen}$. Taking into account the isospin flavour symmetry of strong interactions, we obtain

$$\langle J/\psi K^+ | \mathcal{H}_{\rm eff}^{I=0} | B^+ \rangle = +\langle J/\psi K^0 | \mathcal{H}_{\rm eff}^{I=0} | B_d^0 \rangle \tag{8}$$

$$\langle J/\psi K^+ | \mathcal{H}_{\rm eff}^{I=1} | B^+ \rangle = -\langle J/\psi K^0 | \mathcal{H}_{\rm eff}^{I=1} | B_d^0 \rangle, \tag{9}$$

and finally arrive at

$$A(B^+ \to J/\psi K^+) = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs}V_{cb}^* \left\{ A_c^{(0)} - A_c^{(1)} \right\} + V_{us}V_{ub}^* \left\{ A_u^{(0)} - A_u^{(1)} \right\} \right] \tag{10}$$

$$A(B_d^0 \to J/\psi K^0) = \frac{G_{\rm F}}{\sqrt{2}} \left[V_{cs}V_{cb}^* \left\{ A_c^{(0)} + A_c^{(1)} \right\} + V_{us}V_{ub}^* \left\{ A_u^{(0)} + A_u^{(1)} \right\} \right], \tag{11}$$

where

$$A_c^{(0)} = A_{\rm CC}^{c} - A_{\rm QCD}^{\rm pen} - A_{\rm EW}^{(0)}, \quad A_c^{(1)} = -A_{\rm EW}^{(1)} \tag{12}$$

$$A_u^{(0)} = A_{\rm CC}^{u(0)} - A_{\rm QCD}^{\rm pen} - A_{\rm EW}^{(0)}, \quad A_u^{(1)} = A_{\rm CC}^{u(1)} - A_{\rm EW}^{(1)} \tag{13}$$

denote hadronic matrix elements $\langle J/\psi K|\mathcal{Q}|B\rangle$, i.e. are CP-conserving strong amplitudes. The CKM factors in (10) and (11) are given by

$$V_{cs}V_{cb}^* = \left(1 - \frac{\lambda^2}{2}\right)\lambda^2 A, \quad V_{us}V_{ub}^* = \lambda^4 A\, R_b\, e^{i\gamma}, \tag{14}$$

with

$$\lambda \equiv |V_{us}| = 0.22, \quad A \equiv |V_{cb}|/\lambda^2 = 0.81 \pm 0.06, \quad R_b \equiv |V_{ub}/(\lambda V_{cb})| = 0.41 \pm 0.07, \tag{15}$$

and γ is the usual angle of the unitarity triangle of the CKM matrix [3]. Consequently, we may write

$$A(B^+ \to J/\psi K^+) = \frac{G_{\rm F}}{\sqrt{2}} \left(1 - \frac{\lambda^2}{2}\right) \lambda^2 A \left\{ A_c^{(0)} - A_c^{(1)} \right\} \left[1 + \frac{\lambda^2 R_b}{1-\lambda^2/2} \left\{ \frac{A_u^{(0)} - A_u^{(1)}}{A_c^{(0)} - A_c^{(1)}} \right\} e^{i\gamma} \right] \tag{16}$$

$$A(B_d^0 \to J/\psi K^0) = \frac{G_{\rm F}}{\sqrt{2}} \left(1 - \frac{\lambda^2}{2}\right) \lambda^2 A \left\{ A_c^{(0)} + A_c^{(1)} \right\} \left[1 + \frac{\lambda^2 R_b}{1-\lambda^2/2} \left\{ \frac{A_u^{(0)} + A_u^{(1)}}{A_c^{(0)} + A_c^{(1)}} \right\} e^{i\gamma} \right]. \tag{17}$$

Let us note that (17) takes the same form as the parametrization derived in [13], making, however, its isospin decomposition explicit. An important observation is that the CP-violating phase factor $e^{i\gamma}$ enters in (16) and (17) in a doubly Cabibbo-suppressed way. Moreover, the $A_u^{(0,1)}$ amplitudes are governed by penguin-like topologies and annihilation diagrams, and are hence expected to be suppressed with respect to $A_c^{(0)}$, which originates from tree-diagram-like topologies [14]. In order to keep track of this feature, we introduce a "generic" expansion parameter $\overline{\lambda} = \mathcal{O}(0.2)$ [15], which is of the same order as the Wolfenstein parameter $\lambda = 0.22$:

$$\left| A_u^{(0,1)}/A_c^{(0)} \right| = \mathcal{O}(\overline{\lambda}). \tag{18}$$

Consequently, the $e^{i\gamma}$ terms in (16) and (17) are actually suppressed by $\mathcal{O}(\overline{\lambda}^3)$. Since $A_c^{(1)}$ is due to dynamically suppressed matrix elements of EW penguin operators (see (12)), we expect

$$\left| A_c^{(1)}/A_c^{(0)} \right| = \mathcal{O}(\overline{\lambda}^3). \tag{19}$$

Therefore, we obtain – up to negligibly small corrections of $\mathcal{O}(\overline{\lambda}^3)$ – the following expression:

$$A(B^+ \to J/\psi K^+) = A(B_d^0 \to J/\psi K^0) = A_{\rm SM}^{(0)}, \tag{20}$$

with

$$A_{\rm SM}^{(0)} \equiv \frac{G_{\rm F}}{\sqrt{2}} \left(1 - \frac{\lambda^2}{2}\right) \lambda^2 A\, A_c^{(0)}. \tag{21}$$

Let us note that the plausible hierarchy of strong amplitudes given in (18) and (19) may be spoiled by very large rescattering processes [16]. In the worst case, (20) may receive corrections of $\mathcal{O}(\overline{\lambda}^2)$, and not at the $\overline{\lambda}^3$ level. However, we do not consider this a very likely scenario and note that also the "QCD factorization" approach developed in [11] is not in favour of such large rescattering effects.

The purpose of this paper is to explore the impact of physics beyond the Standard Model on $B \to J/\psi K$ decays. In the next section, we investigate the generic size of such effects using dimensional arguments following from the framework of effective field theory.

3 Effects of Physics Beyond the Standard Model

The generic way of introducing physics beyond the Standard Model is to use the language of effective field theory and to write down all possible dim-6 operators. Of course, this has been known already for a long time and lists of the dim-6 operators involving all the Standard-Model particles have been published in the literature [17]. After having introduced these additional operators, we construct the generalization of the Standard-Model effective Hamiltonian (2) at the scale of the b quark. The relevant operators are again dim-6 operators, the coefficients of which now contain a Standard-Model contribution, and a possible piece of "new physics".

The problem with this point of view is that the list of possible dim-6 operators contains close to one hundred entries, not yet taking into account the flavour structure, making this general approach almost useless. However, we are dealing with non-leptonic decays in which we are, because of our ignorance of the hadronic matrix elements, sensitive neither to the helicity structure of the operators nor to their colour structure. The only information that is relevant in this case is the flavour structure, and hence we introduce the notation

$$\left[(\overline{Q}q_1)(\overline{q}_2 q_3)\right] = \sum [\text{Wilson coefficient}] \times [\text{dim-6 operator mediating } Q \to q_1\overline{q}_2 q_3]. \tag{22}$$

Clearly, this sum is renormalization-group-invariant, and involves, at the scale of the b quark, Standard-Model as well as possible non-Standard-Model contributions. In particular, it allows us to estimate the relative size of a possible new-physics contribution.

3.1 B_d^0–$\overline{B_d^0}$ Mixing: $\Delta B = \pm 2$ Operators

Using this language, we now consider the $\Delta B = \pm 2$ operators relevant to B_d^0–$\overline{B_d^0}$ mixing. We have [18]

$$\mathcal{H}_{\rm eff}(\Delta B = +2) = G \left[(\overline{b}d)(\overline{b}d)\right] \tag{23}$$

as the relevant flavour structure. Within the Standard Model, (23) originates from the well-known box diagrams, which are strongly suppressed by the CKM factor $(V_{td}V_{tb}^*)^2$, as well as by a loop factor

$$\frac{g_2^2}{64\pi^2} = \frac{G_{\rm F} M_W^2}{\sqrt{128}\pi^2} \approx 1 \times 10^{-3}, \tag{24}$$

making the Standard-Model contribution very small, of the order of[1]

$$G_{\rm SM} = \frac{G_{\rm F}}{\sqrt{2}} \left(\frac{G_{\rm F} M_W^2}{\sqrt{128}\pi^2} \right) (V_{td} V_{tb}^*)^2. \tag{25}$$

A new-physics contribution therefore could in principle be as large as the Standard-Model piece. If Λ is the scale of physics beyond the Standard Model, we have

$$G_{\rm NP} = \frac{G_{\rm F}}{\sqrt{2}} \left(\frac{G_{\rm F} M_W^2}{\sqrt{128}\pi^2} \right) \frac{M_W^2}{\Lambda^2} e^{-i2\psi}, \tag{26}$$

where ψ is a possible weak phase, which is induced by the new-physics contribution. Finally, we arrive at

$$G = \frac{G_{\rm F}}{\sqrt{2}} \left(\frac{G_{\rm F} M_W^2}{\sqrt{128}\pi^2} \right) \left[(V_{td} V_{tb}^*)^2 + \frac{M_W^2}{\Lambda^2} e^{-i2\psi} \right] \equiv |R| e^{-i\phi_{\rm M}}. \tag{27}$$

The relevant quantity is the (weak) phase $\phi_{\rm M}$ of this expression, which enters the "mixing-induced" CP-violating asymmetries [3]. Using the standard parametrization

$$V_{td} V_{tb}^* = \lambda^3 A R_t e^{-i\beta}, \tag{28}$$

where $R_t \equiv |V_{td}/(\lambda V_{cb})| = \mathcal{O}(1)$, we obtain

$$\tan \phi_{\rm M} = \frac{\sin(2\beta) + \varrho^2 \sin(2\psi)}{\cos(2\beta) + \varrho^2 \cos(2\psi)}, \tag{29}$$

with

$$\varrho = \left(\frac{1}{\lambda^3 A R_t} \right) \left(\frac{M_W}{\Lambda} \right). \tag{30}$$

Since the factor ϱ can be of order one even for large Λ, there can be a large phase shift in the mixing phase. If we assume, for example, $AR_t = 1$, this term equals 1 for a new-physics scale of $\Lambda \sim 8\,\text{TeV}$. Such contributions affect of course not only the CP-violating phase $\phi_{\rm M}$, but also the "strength" $|R|$ of the B_d^0–$\overline{B_d^0}$ mixing, which would manifest itself as an inconsistency in the usual "standard analysis" of the unitarity triangle [19].

3.2 Decay Amplitudes: $\Delta B = \pm 1$ Operators

We can discuss the operators mediating $\Delta B = \pm 1$ processes on the same footing as B_d^0–$\overline{B_d^0}$ mixing. In the presence of new physics, the corresponding low-energy effective Hamiltonian can also be composed in $I = 0$ and $I = 1$ pieces, as in (7). In Section 2, we had a closer look at the Standard-Model contributions, arising from current–current, QCD and EW penguin operators. New physics may affect the corresponding Wilson coefficients, and may introduce new dim-6 operators, modifying (20) as follows:

$$A(B^+ \to J/\psi K^+) = A_{\rm SM}^{(0)} \left[1 + \sum_k r_0^{(k)} e^{i\delta_0^{(k)}} e^{i\varphi_0^{(k)}} - \sum_j r_1^{(j)} e^{i\delta_1^{(j)}} e^{i\varphi_1^{(j)}} \right] \tag{31}$$

$$A(B_d^0 \to J/\psi K^0) = A_{\rm SM}^{(0)} \left[1 + \sum_k r_0^{(k)} e^{i\delta_0^{(k)}} e^{i\varphi_0^{(k)}} + \sum_j r_1^{(j)} e^{i\delta_1^{(j)}} e^{i\varphi_1^{(j)}} \right]. \tag{32}$$

[1] Note that we do not write the Inami–Lim function coming from the box diagram; this function is of order one and is ignored in our estimates. A similar comment applies to perturbative QCD corrections.

Here $r_0^{(k)}$ and $r_1^{(j)}$ correspond to the $I=0$ and $I=1$ pieces, respectively, $\delta_0^{(k)}$ and $\delta_1^{(j)}$ are CP-conserving strong phases, and $\varphi_0^{(k)}$ and $\varphi_1^{(j)}$ the corresponding CP-violating weak phases. The amplitudes for the CP-conjugate processes can be obtained straightforwardly from (31) and (32) by reversing the signs of the weak phases. The labels k and j distinguish between different new-physics contributions to the $I=0$ and $I=1$ sectors.

For the following discussion, we have to make assumptions about the size of a possible new-physics piece. We shall assume that the new-physics contributions to the $I=0$ sector are smaller compared to the leading Standard-Model amplitude (21) by a factor of order $\overline{\lambda}$, i.e.

$$r_0^{(k)} = \mathcal{O}(\overline{\lambda}). \tag{33}$$

In the case where the new-physics effects are even smaller, it is difficult to disentangle them from the Standard-Model contribution; this will be addressed, together with several other scenarios, in Section 5. Parametrizing the new-physics amplitudes again by a scale Λ, we have

$$\frac{G_{\rm F}}{\sqrt{2}}\frac{M_W^2}{\Lambda^2} \sim \overline{\lambda}\left[\frac{G_{\rm F}}{\sqrt{2}}\,\lambda^2 A\right], \tag{34}$$

corresponding to $\Lambda \sim 1\,\text{TeV}$. Consequently, as in the example given after (30), also here we have a generic new-physics scale in the TeV regime.

As far as possible new-physics contributions to the $I=1$ sector are concerned, we assume a similar "generic strength" of the corresponding operators. However, in comparison with the $I=0$ pieces, the matrix elements of the $I=1$ operators, having the general flavour structure

$$\mathcal{Q} \sim (\overline{u}u - \overline{d}d)(\overline{b}s), \tag{35}$$

are expected to suffer from a dynamical suppression. As in (18) and (19), we shall assume that this brings another factor of $\overline{\lambda}$ into the game, yielding

$$r_1^{(j)} = \mathcal{O}(\overline{\lambda}^2). \tag{36}$$

Employing this kind of counting, the new-physics contributions to the $I=1$ sector would be enhanced by a factor of $\mathcal{O}(\overline{\lambda})$ with respect to the $I=1$ Standard-Model pieces. This may actually be the case if new physics shows up, for example, in EW penguin processes.

Consequently, we obtain

$$A(B\to J/\psi K) = A_{\rm SM}^{(0)}\Big[1 + \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^3)}_{\mathrm{SM}}\Big]. \tag{37}$$

In the presence of large rescattering effects, the assumed dynamical suppression through a factor of $\mathcal{O}(\overline{\lambda})$ would no longer be effective, thereby modifying (37) as follows:

$$A(B\to J/\psi K)|_{\rm res.} = A_{\rm SM}^{(0)}\Big[1 + \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=0}} + \underbrace{\mathcal{O}(\overline{\lambda})}_{\mathrm{NP}_{I=1}} + \underbrace{\mathcal{O}(\overline{\lambda}^2)}_{\mathrm{SM}}\Big]. \tag{38}$$

However, as we have noted above, we do not consider this as a very likely scenario, and shall use (37) in the following discussion, neglecting the Standard-Model pieces of $\mathcal{O}(\overline{\lambda}^3)$, which are not under theoretical control.

Concerning the analysis of CP violation, it is obvious that possible weak phases appearing in the new-physics contributions play the key role. As was the case for the $\Delta B=\pm 2$ operators, also the $\Delta B=\pm 1$ operators could carry such new weak phases, which would then affect the CP-violating $B\to J/\psi K$ observables.

4 Observables of $B \to J/\psi K$ Decays

The decays $B^+ \to J/\psi K^+$, $B_d^0 \to J/\psi K^0$ and their charge conjugates provide a set of four decay amplitudes A_i. Measuring the corresponding rates, we may determine the $|A_i|^2$. Since we are not interested in the overall normalization of the decay amplitudes, we may construct three independent observables with the help of the $|A_i|^2$:

$$\mathcal{A}_{\rm CP}^{(+)} \equiv \frac{|A(B^+ \to J/\psi K^+)|^2 - |A(B^- \to J/\psi K^-)|^2}{|A(B^+ \to J/\psi K^+)|^2 + |A(B^- \to J/\psi K^-)|^2} \tag{39}$$

$$\mathcal{A}_{\rm CP}^{\rm dir} \equiv \frac{|A(B_d^0 \to J/\psi K^0)|^2 - |A(\overline{B_d^0} \to J/\psi \overline{K^0})|^2}{|A(B_d^0 \to J/\psi K^0)|^2 + |A(\overline{B_d^0} \to J/\psi \overline{K^0})|^2} \tag{40}$$

$$B \equiv \frac{\langle |A(B_d \to J/\psi K)|^2\rangle - \langle |A(B^\pm \to J/\psi K^\pm)|^2\rangle}{\langle |A(B_d \to J/\psi K)|^2\rangle + \langle |A(B^\pm \to J/\psi K^\pm)|^2\rangle}, \tag{41}$$

where the "CP-averaged" amplitudes are defined as follows:

$$\langle |A(B_d \to J/\psi K)|^2\rangle \equiv \frac{1}{2}\left[|A(B_d^0 \to J/\psi K^0)|^2 + |A(\overline{B_d^0} \to J/\psi \overline{K^0})|^2\right] \tag{42}$$

$$\langle |A(B^\pm \to J/\psi K^\pm)|^2\rangle \equiv \frac{1}{2}\left[|A(B^+ \to J/\psi K^+)|^2 + |A(B^- \to J/\psi K^-)|^2\right]. \tag{43}$$

If we consider the neutral decays $B_d \to J/\psi K_{\rm S}$, where the final state is a CP eigenstate with eigenvalue -1, the following time-dependent CP asymmetry provides an additional observable, which is due to interference between B_d^0–$\overline{B_d^0}$ mixing and decay processes [3]:

$$\frac{\Gamma(B_d^0(t) \to J/\psi K_{\rm S}) - \Gamma(\overline{B_d^0}(t) \to J/\psi K_{\rm S})}{\Gamma(B_d^0(t) \to J/\psi K_{\rm S}) + \Gamma(\overline{B_d^0}(t) \to J/\psi K_{\rm S})} = \mathcal{A}_{\rm CP}^{\rm dir}\,\cos(\Delta M_d t) + \mathcal{A}_{\rm CP}^{\rm mix}\,\sin(\Delta M_d t). \tag{44}$$

Here the rates correspond to decays of initially, i.e. at time $t = 0$, present B_d^0- or $\overline{B_d^0}$-mesons, and ΔM_d denotes the mass difference between the B_d mass eigenstates. The "direct" CP-violating contribution $\mathcal{A}_{\rm CP}^{\rm dir}$ was already introduced in (40), and the "mixing-induced" CP asymmetry is given by

$$\mathcal{A}_{\rm CP}^{\rm mix} = \frac{2\,{\rm Im}\,\xi}{1+|\xi|^2}\,, \tag{45}$$

with

$$\xi = e^{-i\phi}\left[\frac{1+\sum_k r_0^{(k)} e^{i\delta_0^{(k)}}\,e^{-i\varphi_0^{(k)}} + \sum_j r_1^{(j)} e^{i\delta_1^{(j)}}\,e^{-i\varphi_1^{(j)}}}{1+\sum_k r_0^{(k)} e^{i\delta_0^{(k)}}\,e^{+i\varphi_0^{(k)}} + \sum_j r_1^{(j)} e^{i\delta_1^{(j)}}\,e^{+i\varphi_1^{(j)}}}\right]. \tag{46}$$

In (46), we have used the parametrization (32) to express the corresponding decay amplitudes. The CP-violating weak phase ϕ is given by $\phi = \phi_{\rm M} + \phi_K$, where $\phi_{\rm M}$ was introduced in (27), and ϕ_K is the weak K^0–$\overline{K^0}$ mixing phase, which is negligibly small in the Standard Model. Owing to the small value of the CP-violating parameter ε_K of the neutral kaon system, ϕ_K can only be affected by very contrived models of new physics [20].

In order to search for new-physics effects in the $B \to J/\psi K$ system, it is useful to introduce the following combinations of the observables (39) and (40):

$$S \equiv \frac{1}{2}\left[\mathcal{A}_{\rm CP}^{\rm dir} + \mathcal{A}_{\rm CP}^{(+)}\right], \quad D \equiv \frac{1}{2}\left[\mathcal{A}_{\rm CP}^{\rm dir} - \mathcal{A}_{\rm CP}^{(+)}\right]. \tag{47}$$

Using the parametrizations (31) and (32), and assuming the hierarchy in (37), we obtain

$$S = -2\left[\sum_k r_0^{(k)} \sin\delta_0^{(k)} \sin\varphi_0^{(k)}\right]\left[1 - 2\sum_l r_0^{(l)} \cos\delta_0^{(l)} \cos\varphi_0^{(l)}\right] = \mathcal{O}(\overline{\lambda}) + \mathcal{O}(\overline{\lambda}^2) \tag{48}$$

$$D = -2\sum_j r_1^{(j)} \sin\delta_1^{(j)} \sin\varphi_1^{(j)} = \mathcal{O}(\overline{\lambda}^2) \tag{49}$$

$$B = +2\sum_j r_1^{(j)} \cos\delta_1^{(j)} \cos\varphi_1^{(j)} = \mathcal{O}(\overline{\lambda}^2), \tag{50}$$

where terms of $\mathcal{O}(\overline{\lambda}^3)$, including also a Standard-Model contribution, which is not under theoretical control, have been neglected. Note that if the dynamical suppression of the $I = 1$ contributions would be larger, B and D would be further suppressed relative to S.

The corresponding expression for the mixing-induced CP asymmetry (45) is rather complicated and not very instructive. Let us give it for the special case where the new-physics contributions to the $I = 0$ and $I = 1$ sectors involve either the same weak or strong phases:

$$\begin{aligned}\mathcal{A}_{\rm CP}^{\rm mix} = &-\sin\phi - 2\,r_0\cos\delta_0\sin\varphi_0\cos\phi - 2\,r_1\cos\delta_1\sin\varphi_1\cos\phi \\ &+ r_0^2\Big[(1-\cos 2\varphi_0)\sin\phi + \cos 2\delta_0\sin 2\varphi_0\cos\phi\Big] = -\sin\phi + \mathcal{O}(\overline{\lambda}) + \mathcal{O}(\overline{\lambda}^2).\end{aligned} \tag{51}$$

Expressions (48)–(50) also simplify in this case:

$$S = -2\,r_0\sin\delta_0\sin\varphi_0 + r_0^2\sin 2\delta_0\sin 2\varphi_0,\; D = -2\,r_1\sin\delta_1\sin\varphi_1,\; B = 2\,r_1\cos\delta_1\cos\varphi_1. \tag{52}$$

In the following section, we discuss the search for new physics with these observables in more detail, and have also a closer look at the present experimental situation.

5 Discussion

Let us begin our discussion by turning first to the present experimental situation. Concerning the direct CP asymmetries (39) and (40), we have

$$\mathcal{A}_{\rm CP}^{\rm dir} = (26 \pm 19)\,\% \;\; \text{(BaBar [6])}, \quad \mathcal{A}_{\rm CP}^{(+)} = \begin{cases}(13 \pm 14)\% & \text{(BaBar [6])}\\ (-1.8 \pm 4.3 \pm 0.4)\% & \text{(CLEO [21]).}\end{cases} \tag{53}$$

The present status of the mixing-induced $B_d \to J/\psi K_{\rm S}$ CP asymmetry (45) is given as follows:[2]

$$-\mathcal{A}_{\rm CP}^{\rm mix} = \begin{cases} 0.79^{+0.41}_{-0.44} & \text{(CDF [5])}\\ -0.10 \pm 0.42 & \text{(BaBar [6])}\\ 0.49^{+0.53}_{-0.57} & \text{(Belle [7]).}\end{cases} \tag{54}$$

[2]The values for $\sin 2\beta$ reported in [6, 7] are dominated by $B_d \to J/\psi K_{\rm S}$, but actually correspond to an average over various modes.

Using the rates listed in [22], as well as $\tau_{B^+}/\tau_{B^0_d} = 1.060 \pm 0.029$, and adding the experimental errors in quadrature, we obtain[3]

$$B = \frac{\mathrm{BR}(B_d \to J/\psi K)\tau_{B^+}/\tau_{B^0_d} - \mathrm{BR}(B^\pm \to J/\psi K^\pm)}{\mathrm{BR}(B_d \to J/\psi K)\tau_{B^+}/\tau_{B^0_d} + \mathrm{BR}(B^\pm \to J/\psi K^\pm)} = (-2.9 \pm 4.3)\%, \tag{55}$$

where the numerical value depends rather sensitively on the lifetime ratio. On the other hand, the BaBar results for the CP asymmetries listed in (53) yield

$$S = (20 \pm 12)\%, \quad D = (7 \pm 12)\%. \tag{56}$$

In view of the large experimental uncertainties, we cannot yet draw any conclusions. However, the situation should improve significantly in the future. As can be seen in (48)–(50), the observable S provides a "smoking-gun" signal of new-physics contributions to the $I = 0$ sector, while D and B allow us to probe new physics affecting the $I = 1$ pieces. As the hierarchy in (37) implies that S receives terms of $\mathcal{O}(\overline{\lambda})$, whereas D and B arise both at the $\overline{\lambda}^2$ level, we conclude that S may already be accessible at the first-generation B-factories (BaBar, Belle, Tevatron-II), whereas the latter observables will probably be left for second-generation B experiments (BTeV, LHCb). However, should B and D, in addition to S, also be found to be at the 10% level, i.e. should be measured at the first-generation B-factories, we would not only have signals for physics beyond the Standard Model, but also for large rescattering processes.

A more pessimistic scenario one can imagine is that S is measured at the $\overline{\lambda}^2$ level in the LHC era, whereas no indications for non-vanishing values of D and B are found. Then we would still have evidence for new physics, which would then correspond to $r_0^{(k)} = \mathcal{O}(\overline{\lambda}^2)$ and $r_1^{(j)} = \mathcal{O}(\overline{\lambda}^3)$. However, if all three observables are measured to be of $\mathcal{O}(\overline{\lambda}^2)$, new-physics effects cannot be distinguished from Standard-Model contributions, which could also be enhanced to the $\overline{\lambda}^2$ level by large rescattering effects. This would be the most unfortunate case for the strategy to search for new-physics contributions to the $B \to J/\psi K$ decay amplitudes proposed in this paper. However, further information can be obtained with the help of the decay $B_s \to J/\psi K_{\mathrm{S}}$, which can be combined with $B_d \to J/\psi K_{\mathrm{S}}$ through the U-spin symmetry of strong interactions to extract the angle γ of the unitarity triangle, and may shed light on new physics even in this case [13].

As can be seen in (51), the mixing-induced CP asymmetry is affected both by $I = 0$ and by $I = 1$ new-physics contributions, where the dominant $\mathcal{O}(\overline{\lambda})$ effects are expected to be due to the $I = 0$ sector. Neglecting terms of $\mathcal{O}(\overline{\lambda}^2)$, we may write

$$\mathcal{A}_{\mathrm{CP}}^{\mathrm{mix}} = -\sin(\phi + \delta\phi_{\mathrm{NP}}^{\mathrm{dir}}), \tag{57}$$

with

$$\delta\phi_{\mathrm{NP}}^{\mathrm{dir}} = 2\sum_k r_0^{(k)} \cos\delta_0^{(k)} \sin\varphi_0^{(k)}. \tag{58}$$

The phase shift $\delta\phi_{\mathrm{NP}}^{\mathrm{dir}} = \mathcal{O}(\overline{\lambda})$ may be as large as $\mathcal{O}(20^\circ)$. Since the Standard-Model range for ϕ is given by $28^\circ \leq \phi = 2\beta \leq 70^\circ$ [19], the mixing-induced CP asymmetry (57) may also be

[3] Phase-space effects play a negligible role in this expression.

affected significantly by new-physics contributions to the $B_d \to J/\psi K_{\rm S}$ decay amplitude, and not only in the "standard" fashion, through B_d^0–$\overline{B_d^0}$ mixing, as discussed in Subsection 3.1. This would be another possibility to accommodate the small central value of the BaBar result in (54), which was not pointed out in [8]. In order to gain confidence in such a scenario, it is crucial to improve also the measurements of the observables S, D and B. Let us note that if, in addition to S, also D and B should be found to be sizeable, i.e. of $\mathcal{O}(\overline{\lambda})$, also the terms linear in $r_1^{(j)}$ have to be included in (58):

$$\delta\phi_{\rm NP}^{\rm dir}\Big|_{\rm res.} = 2\left[\sum_k r_0^{(k)}\cos\delta_0^{(k)}\sin\varphi_0^{(k)} + \sum_j r_1^{(j)}\cos\delta_1^{(j)}\sin\varphi_1^{(j)}\right]. \tag{59}$$

So far, our considerations were completely general. Let us therefore comment briefly on a special case, where the strong phases $\delta_0^{(k)}$ and $\delta_1^{(j)}$ take the trivial values 0 or π, as in "factorization". In this case, (48)–(50) would simplify as follows:

$$S \approx 0, \quad D \approx 0, \quad B \approx 2\sum_j r_1^{(j)}\sin\varphi_1^{(j)} = \mathcal{O}(\overline{\lambda}^2), \tag{60}$$

whereas (51) would yield

$$\begin{aligned}\mathcal{A}_{\rm CP}^{\rm mix} &= -\sin\phi - 2\,r_0\sin\varphi_0\cos\phi - 2\,r_1\sin\varphi_1\cos\phi \\ &\quad + r_0^2\Big[(1-\cos 2\varphi_0)\sin\phi + \sin 2\varphi_0\cos\phi\Big] = -\sin\phi + \mathcal{O}(\overline{\lambda}) + \mathcal{O}(\overline{\lambda}^2).\end{aligned} \tag{61}$$

The important point is that S and D are governed by sines of the strong phases, whereas the new-physics contributions to B and $\mathcal{A}_{\rm CP}^{\rm mix}$ involve cosines of the corresponding strong phases. Consequently, these terms do *not* vanish for $\delta \to 0, \pi$. The impact of new physics on $\mathcal{A}_{\rm CP}^{\rm mix}$ may still be sizeable in this scenario, whereas B could only be measured in the LHC era [23]. If S and D should be observed at the $\overline{\lambda}$ and $\overline{\lambda}^2$ levels, respectively, we would not only get a "smoking-gun" signal for new-physics contributions to the $B \to J/\psi K$ decay amplitudes, but also for non-factorizable hadronic effects. A measurement of all three observables S, D and B at the $\overline{\lambda}$ level would imply, in addition, large rescattering processes, as we have already emphasized above.

6 Summary

We have presented a general analysis of new-physics effects in the $B^\pm \to J/\psi K^\pm$, $B_d \to J/\psi K_{\rm S}$ system. To this end, we have taken into account the constraints that are implied by the $SU(2)$ isospin symmetry of strong interactions, and have estimated the generic size of the new-physics contributions through dimensional arguments following from the picture of effective field theory. In addition to the usual mixing-induced CP asymmetry $\mathcal{A}_{\rm CP}^{\rm mix}$ of the $B_d \to J/\psi K_{\rm S}$ mode, we have introduced a set of three observables, S, D and B, which play the key role to search for "smoking-gun" signals of new-physics contributions to the $I = 0$ and $I = 1$ decay-amplitude sectors. Imposing a plausible dynamical hierarchy of amplitudes, we argue that S may already be accessible at the first-generation B-factories, whereas the remaining

ones will probably be left for second-generation B experiments of the LHC era. However, in the presence of large rescattering effects, all three new-physics observables S, D and B may be sizeable. At present, the large experimental uncertainties on these quantities do not allow us to draw any conclusions, and we strongly encourage our experimental colleagues to focus not only on $\mathcal{A}_{\rm CP}^{\rm mix}$, but also on S, D and B. A future measurement corresponding to the central values given in (56), i.e. $S = 20\%$ and $D = 7\%$, would be as exciting as $\mathcal{A}_{\rm CP}^{\rm mix} = 10\%$. Also the latter result could be due to new-physics contributions to the $B_d \to J/\psi K_{\rm S}$ decay amplitudes, and would not necessarily be an indication of new physics in B_d^0–$\overline{B_d^0}$ mixing. We look forward to better data on $B \to J/\psi K$ decays, which will, hopefully, open a window to the physics beyond the Standard Model.

Acknowledgements

The work of T.M. is supported by the DFG Graduiertenkolleg "Elementarteilchenphysik an Beschleunigern", by the DFG Forschergruppe "Quantenfeldtheorie, Computeralgebra und Monte Carlo Simulationen", and by the "Ministerium für Bildung und Forschung" (bmb+f).

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